Asymmetric Sneutrino Dark Matter and the $\Omega_{\mathrm{b}}/\Omega_{\mathrm{DM}}$ Puzzle

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Abstract

The inferred values of the cosmological baryon and dark matter densities are strikingly similar, but in most theories of the early universe there is no true explanation of this fact; in particular, the baryon asymmetry and thus density depends upon unknown, and *a priori* small, CP-violating phases which are independent of all parameters determining the dark matter density. We consider models of dark matter possessing a particle-antiparticle asymmetry where this asymmetry determines both the baryon asymmetry and strongly effects the dark matter density, thus naturally linking Ω_{b} and Ω_{dm}. We show that sneutrinos can play the role of such dark matter in a previously studied variant of the MSSM in which the light neutrino masses result from higher-dimensional supersymmetry-breaking terms.

1   Introduction

For some time it has been apparent that the inferred values of the cosmological baryon and dark matter densities are strikingly similar. For example, the latest WMAP-determined range for the dark matter density, $0.129 > \Omega_{\mathrm{CDM}}h^2 > 0.095$, is within a factor of a few of the combined WMAP and big-bang nucleosynthesis determined value of the baryon density, $0.025 > \Omega_{\mathrm{b}}h^2 > 0.012$ [1, 2]. In the vast majority of models of the early universe, the cosmological baryon and dark matter densities are independently determined. The surviving baryon density is set by a baryon asymmetry generated during baryogenesis, and thus depends upon unknown, and *a priori* small CP-violating phases, as well as unknown baryon-number violating dynamics. In contrast, the dark matter density is set by the 'freeze-out' of the interactions that keep the dark matter in equilibrium, and is independent of the dynamics of baryogenesis. In particular, although weakly-interacting massive particles at the TeV-scale naturally have a relic density of $\mathcal{O}(1)$ times the critical density, this is not at all the case for baryons. The Boltzmann equations show the size of the baryon relic density in the absence of an asymmetry is $\mathcal{O}(10^{-11})$. Thus the comparative closeness of the baryon and dark matter densities poses a puzzle. One possible approach to this problem is to more closely integrate the dynamics of baryogenesis with that of the origin of dark matter. In particular, it is natural to consider models where the dark matter and baryon sectors share a quantum number, either continuous or discrete, which provides a relation between their surviving number densities and thus energy densities.[1]

Specifically, in this letter we consider models of dark matter possessing a particle-antiparticle asymmetry where this asymmetry strongly effects the dark matter density, and through the electroweak (EW) anomaly, determines the baryon asymmetry, thus naturally linking Ω_{b} and Ω_{CDM}. (A noteworthy early attempt along these lines which shares some features with our model is contained in Ref. [5].) The shared quantum number clearly requires that the dark matter particle not be it's own antiparticle, so the putative dark matter candidate cannot be the standard LSP neutralino of the broken Minimal Supersymmetric Standard Model (MSSM). In fact, we will show in sections 3 and 4 that sneutrinos can play the role of such dark matter in a previously studied variant of the MSSM in which the light neutrino masses result from higher-dimensional supersymmetry-breaking terms [6, 7, 8, 9, 10]. This model preserves all the successes of the MSSM, such as stability of the weak scale and unification of gauge couplings, while being, at least in part, testable at the LHC.

Before we proceed to the calculation of the relic dark matter and baryon densities and the details of our model, it is useful to consider some aspects of the idea of a shared quantum number determining the ratio of baryon to dark matter. First, consider the simplified case in which, to a very good approximation, the dark matter and baryon sectors cannot exchange their conserved quantum number after its first production. As an example, consider the situation in which the asymptotic baryon number $B = 1$ states of the Standard Model (SM) have global charge, q, while the lightest asymptotic states

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[1]An alternate approach is to invoke a form of the Anthropic Principle, see, for example, Ref. [3]. Another recent idea is explored in Ref. [4].

(with mass m_{dm}) in a hidden sector carry charge, Q. Then conservation of global charge implies that $q\left(n_{\mathrm{b}} - n_{\overline{\mathrm{b}}}\right) = -Q\left(n_{\mathrm{dm}} - n_{\overline{\mathrm{dm}}}\right)$, where the n's are the number densities of the indicated (anti)particles. Further assume that interactions in both sectors are strong enough such that almost all antiparticles are eliminated by annihilation with their particles (assuming $Q/q < 0$ for simplicity). This implies

$$n_{\mathrm{b}} = cn_{\mathrm{dm}} \quad \text{with} \quad c = |Q/q|, \tag{1}$$

which in turn leads to

$$\frac{\Omega_{\mathrm{b}}}{\Omega_{\mathrm{dm}}} = \frac{m_{\mathrm{b}}}{m_{\mathrm{dm}}} \frac{n_{\mathrm{b}}}{n_{\mathrm{dm}}} = c\frac{m_{\mathrm{b}}}{m_{\mathrm{dm}}}. \tag{2}$$

Thus the energy densities are indeed related, but the ratio is only naturally $\mathcal{O}(1)$ if the ratio of the baryon to dark matter masses is not too small. It is important to note that this disfavours models of the above type where the dark matter candidate is a 'hidden sector' particle and favours a particular class of models where the dark matter candidate has weak scale mass, its mass arising either from electroweak symmetry breaking, or from the dynamics which drives electroweak breaking, such as supersymmetry softly broken at the weak-scale.[2]

As a consequence, the naive statement of Eq.(1) also requires modification as the approximation of negligible interactions between the two sectors, and thus negligible exchange of the shared conserved quantum number, does not hold in realistic models. In this situation, the ratio of the conserved quantum number stored in the two sectors, c, is not simply given by Q/q, but is determined by the 'chemical' equilibrium conditions between the two sectors just before the freeze-out of the relevant interactions. We will further discuss this point in section 2. Moreover, two other alterations to the naive relation Eq.(2) can be present. First, the shared quantum number may only be conserved modulo n, in which case 'self' annihilations may occur, and second, it is not always the case that dark matter interactions are efficient enough such that $n_{\overline{\mathrm{dm}}} \ll n_{\mathrm{dm}}$. Nevertheless, as we argue in the next section, a careful consideration of such models shows that a version of the naive statement Eq.(2) can indeed hold.

2 The Effect of a Matter-Antimatter Asymmetry on the Density of a Thermal Relic

If there were an asymmetry between the density of dark matter particles and dark antimatter particles in the early universe, it may have a considerable effect on the relic density calculation [11]. In this section, we review the relic density calculation for a single species, two species with no asymmetry and two species with an asymmetry.

If dark matter consists of a single species, the relic density calculation is straightforward. The most common example of a single species dark matter candidate would be

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[2]In extensions of the SM such as the MSSM where many gauge-non-neutral fermionic and scalar states also gain mass at this scale, the beta-function coefficient for $SU(3)_{\mathrm{QCD}}$ sharply increases as energies are reduced below the weak scale, and the confinement scale and thus baryon mass is naturally not very much smaller than m_{weak}.

a neutralino LSP with negligible coannihilations with other SUSY particles. In such a case, the resulting relic density is simply

$$\Omega h^2 \cong \frac{1.04 \times 10^9}{M_{\mathrm{Pl}}} \frac{x_F}{\sqrt{g_*}} \frac{1}{(a + 3b/x_F)}, \tag{3}$$

where M_{Pl} is the Planck mass, x_F is the WIMP's mass over the freeze-out temperature, g_* is number of relativistic degrees of freedom available at freeze-out and a and b are related to the WIMP's self-annihilation cross section by $\sigma v = a + bv^2 + \mathcal{O}(v^4)$, v being the relative velocity of the annihilating WIMPs. For a weakly interacting relic, x_F is generally about 20 and varies only mildly with changes in the cross section or mass.

If we instead consider a dark matter candidate *and* its antiparticle, the calculation becomes somewhat different. Labeling these species 1 and 2, we now have each particle's self-annihilation cross section, $\sigma_{11} = \sigma_{22}$, and the particle-antiparticle annihilation cross section, $\sigma_{12} = \sigma_{21}$. In the limit that σ_{12} is much smaller than σ_{11}, we simply get two decoupled species which each provide the relic density of Eq.(3). Otherwise, the effective cross section becomes $\sigma_{11} + \sigma_{12}$ for each species. For the two species together, a density of

$$\Omega h^2 \cong 2 \times \frac{1.04 \times 10^9}{M_{\mathrm{Pl}}} \frac{x_F}{\sqrt{g_*}} \frac{1}{(a_{11} + a_{12} + 3(b_{11} + b_{12})/x_F)} \tag{4}$$

is produced. This case is essentially the calculation of the density of a relic with the coannihilation of a degenerate particle [12].

If an asymmetry exists between the number density of particles and antiparticles in the early universe, again this calculation becomes modified. The most well known example of this is the baryon-antibaryon asymmetry which leads to the current baryon density of the universe despite the very large baryon-antibaryon annihilation cross section. In such a case, where the particle-particle and antiparitcle-antiparticle annihilation cross sections (σ_{11} and σ_{22}) are zero, or negligibly small, the minimum relic density can be related to the magnitude of the particle-antiparticle asymmetry:

$$\Omega h^2_{\min} = \Omega h^2_{\mathrm{bary}} \frac{A}{A_{\mathrm{bary}}} \frac{m}{m_{\mathrm{bary}}}, \tag{5}$$

where $\Omega h^2_{\mathrm{bary}}$ is the current density of baryons in the universe, A and A_{bary} are the particle-antiparticle asymmetries of our relic and of baryons, defined by $A = (n - \overline{n})/n$. A_{bary} is known to be of the order of 10^{-9}. m and m_{bary} are the masses of our relic and of baryons (*i.e.* the proton mass). Of course, if the particle-particle annihilation cross section for the relic is not negligible, Eq.(5) will not hold. There will be a generic tendency for the density of a relic to move towards this value as a result of an asymmetry, however.

To accurately assess the effect of an asymmetry on the thermal relic abundance, the effective cross sections must be determined by integrating over the thermal history. In particular, for each species separately, the densities predicted are

$$\Omega_1 h^2 \cong \frac{1.04 \times 10^9}{M_{\mathrm{Pl}}} \frac{x_F}{\sqrt{g_*}} \frac{1}{(a_{11,\mathrm{eff}} + a_{12,\mathrm{eff}} + 3(b_{11,\mathrm{eff}} + b_{12,\mathrm{eff}})/x_F)} \tag{6}$$

and

$$\Omega_2 h^2 \cong \frac{1.04 \times 10^9}{M_{\mathrm{Pl}}} \frac{x_F}{\sqrt{g_*}} \frac{1}{(a_{22,\mathrm{eff}} + a_{21,\mathrm{eff}} + 3(b_{22,\mathrm{eff}} + b_{21,\mathrm{eff}})/x_F)}, \qquad (7)$$

where the effective cross sections are given by

$$
\begin{aligned}
a_{11,\mathrm{eff}} &= a_{11}, & a_{12,\mathrm{eff}} &= x_F \int_{x_F}^{\infty} \frac{a_{12}}{x^2} \left(\frac{e^{-x}}{e^{-x}+A} \right) dx, \\
a_{22,\mathrm{eff}} &= a_{22}, & a_{21,\mathrm{eff}} &= x_F \int_{x_F}^{\infty} \frac{a_{12}}{x^2} \left(\frac{e^{-x}+A}{e^{-x}} \right) dx, \\
b_{11,\mathrm{eff}} &= b_{11}, & b_{12,\mathrm{eff}} &= 2x_F^2 \int_{x_F}^{\infty} \frac{b_{12}}{x^3} \left(\frac{e^{-x}}{e^{-x}+A} \right) dx, \\
b_{22,\mathrm{eff}} &= b_{22}, & b_{21,\mathrm{eff}} &= 2x_F^2 \int_{x_F}^{\infty} \frac{b_{12}}{x^3} \left(\frac{e^{-x}+A}{e^{-x}} \right) dx.
\end{aligned}
\qquad (8)
$$

The factor, $e^{-x}/(e^{-x} + A)$, is the depletion of the second species due to the asymmetry at a temperature, $T = m/x$. The effect of this depletion depends on the size of the asymmetry and the cross sections involved. If the estimate of Eq.(5) is considerably less than the standard two-species result of Eq.(4), then the asymmetry has little impact and the sum of Eqs.(6) and (7) is nearly equal to Eq.(4). If, however, the estimate of Eq.(5) is larger than the standard two-species result, the relic density moves from the value found with Eq.(4) towards that of Eq.(5). If σ_{11}/σ_{12} is small, then the result will be closer to Eq.(5). If σ_{11}/σ_{12} is large, then the density will be less modified. Note that in some cases, the effect of an asymmetry can be to lower the density of a relic by up to a factor of two. This occurs when σ_{11} and σ_{12} are comparable thus efficiently deplete the second species while leaving the first largely uneffected.

Although these effects are rather model dependent, the net effect can be large regions of parameter space which predict the observed density of dark matter for a WIMP with a mass near $m \sim m_p \Omega_{\mathrm{WMAP}} h^2 / \Omega_{\mathrm{bary}} h^2$, for a matter-antimatter asymmetry similar to that for baryons. For concreteness, we consider a specific example of such a model in the following section.

3   An Example: Mixed Sneutrino Dark Matter

Within the context of the Minimal Supersymmetric Standard Model (MSSM), sneutrinos do not make a very appealing dark matter candidate. They annihilate quite efficiently, resulting in a relic density smaller than the observed dark matter density for sneutrinos lighter than about 500 GeV. Furthermore, their elastic scattering cross section is sufficiently large to be easily observed by direct dark matter experiments. Thus it is natural that in contrast to other supersymmetric dark matter candidates, sneutrino dark matter has received considerably less attention [13].

Looking beyond the vanilla MSSM, however, it is quite easy to construct a viable model for sneutrino dark matter. For example, sneutrinos could mix through A-terms with a singlet forming a mixed sneutrino state. This would reduce the annihilation cross section, potentially providing the appropriate quantity of dark matter, as well as reduce the elastic scattering cross section to experimentally acceptable levels. In fact, such a

scenario has previously been discussed as a way to reconcile the results of the direct dark matter detection experiments DAMA and CDMS [8]. These models have light sneutrino mass eigenstates which are mixtures of left-handed 'active' sneutrinos, and right-handed 'sterile' sneutrinos.

Such mixing naturally arises in models in which the light neutrino masses are generated by supersymmetry breaking effects [6, 7, 9, 10]. The importance to us of these models is that the light sneutrino states share a non-anomalous $(B-L)$-symmetry with the baryons which is only weakly broken by the Majorana neutrino masses. In analogy to the Giudice-Masiero mechanism for the Higgs μ-term there exists a global symmetry that forbids a direct mass, $M_R NN$, and neutrino Yukawa coupling, $\lambda L N H_u$, involving the SM singlet right-handed neutrino superfield N, but allows non-renormalizable couplings to a supersymmetry-breaking hidden-sector field T. The right-handed neutrino mass and the neutrino Yukawa coupling, $\lambda L N H_u$, then only arise from $1/M_{\text{planck}}$-suppressed terms involving the fundamental supersymmetry breaking scale $m_I \sim (v M_{pl})^{1/2}$, where v is the weak scale. The relevant terms in the Lagrangian are

$$\mathcal{L} = \int d^4\theta \left(h \frac{T^\dagger}{M} NN + \tilde{h} \frac{T^\dagger T}{M^2} N^\dagger N + h_B \frac{T^\dagger T T^\dagger}{M^3} NN + \dots \right) + \int d^2\theta \left(g \frac{T}{M} L N H_u \right), \quad (9)$$

where all dimensionless couplings g, h, etc, are taken to be $\mathcal{O}(1)$. The structure of the Lagrangian can be enforced by some simple symmetries [6, 7, 9], including R-parity, which leads to a stable lightest supersymmetric particle (LSP) dark matter candidate. After supersymmetry is broken in the hidden sector at the scale m_I, the F and A components of the field T acquire intermediate scale vevs, $\langle F_T \rangle \sim \langle A_T \rangle^2 \sim m_I^2$. Substituting these vevs into the Lagrangian shows that after SUSY breaking, there arises: (1) a neutrino Yukawa coupling of size $\sim 10^{-7}$–10^{-8}, (2) right-handed neutrino and sneutrino masses at the *weak scale*, (3) a weak-scale trilinear scalar A-term, and (4) a right-handed sneutrino lepton-number violating B-term with magnitude $B^2 \sim (\text{few} \times 100\,\text{MeV})^2$. These together lead to two sources of neutrino masses, a tree level see-saw contribution, as well as a dominant 1-loop contribution, leading to a consistent and attractive neutrino phenomenology [6, 9].

More relevant for us, the sneutrino mass matrix squared is

$$\begin{pmatrix} \tilde{\nu}^* & n \end{pmatrix} \begin{pmatrix} m_L^2 & \frac{1}{\sqrt{2}} A v \sin\beta \\ \frac{1}{\sqrt{2}} A v \sin\beta & m_R^2 \end{pmatrix} \begin{pmatrix} \tilde{\nu} \\ n^* \end{pmatrix}, \quad (10)$$

where we have ignored the small lepton number violating B-term for the moment. Therefore mixing between the left-handed 'active' sneutrino, $\tilde{\nu}$, and the right-handed 'sterile' sneutrino, $\tilde{n}$, occurs, producing the light mass eigenstate $\tilde{\nu}_1 = -\tilde{\nu} \sin\theta + \tilde{n}^* \cos\theta$, where $\tan 2\theta = -\sqrt{2} A v \sin\beta/(m_R^2 - m_L^2)$, with $m_R^2 = M_N^2 + m_{R,\text{soft}}^2$ and $m_L^2 = m_Z^2 \cos 2\beta/2 + m_{L,\text{soft}}^2$. Since the coupling of the lighter sneutrino eigenstate to the Z is suppressed by $\sin\theta$, the direct LEP experimental constraints are weakened, and the calculation of the sneutrino relic abundances is modified. A numerical analysis of the susy parameter space shows that a mixing $\sin\theta \simeq 0.3$ is quite natural.[3] (The B^2 mass-term splits the CP-even and CP-odd components of $\tilde{\nu}_1$ by the amount $\delta \sim \cos^2\theta B^2/m_1$.

[3]Note that the size of the A terms is limited by considerations of vacuum stability [10].

This can have important consequences for detection, potentially making consistent the DAMA and CDMS experiments depending upon parameters [8]. For the purposes of our discussion this splitting is irrelevant.)

We now turn to the calculation of the relative asymmetry in the sneutrino and baryon sectors. The method is a simple adaption of the standard 'chemical' equilibriation techniques applied to, for example, the calculation of the ratio $B/(B-L)$ in the SM [15] or MSSM [16] in the presence of anomaly-induced baryon number violating processes in the early universe. For each species of particle we introduce a chemical potential, μ_A, and for each interaction in equilibrium, say $A + B \leftrightarrow C + D$ there is a corresponding relation among chemical potentials $\mu_A + \mu_B = \mu_C + \mu_D$. Note that when in *thermal* (not necessarily chemical) equilibrium the CP-conjugate species $\overline{A}$ has $\mu_{\overline{A}} = -\mu_A$. The asymmetry in particle-antiparticle number density for the relevant case of $\mu \ll T$ and a species of mass m is given by $n - \overline{n} = gT^2\mu/6$ for $T \gg m$, and is Boltzmann suppressed by $\exp(-m/T)$ for $T \lesssim m$.

The electroweak $B+L$ anomaly interactions provide an efficient mechanism to transmute between B- and L-number densities, at least above the EW phase transition temperature T_c where they are rapid. Our task is to calculate the ratio of sneutrino to baryon asymmetries at T_c where the exchange between the two sectors switches off. (We implicitly assume that the EWPT is 1st order so that anomaly-induced interactions drop out of equilibrium below T_c. For realistic values of T_c the results of our calculation are not particularly sensitive to this assumption.) Above the EWPT temperature the chemical potentials for all SM gauge bosons vanish, so, for example, members of an EW doublet have the same μ. Moreover, efficient flavour mixing among quarks induced by EW interactions such as $u_{Li} + \phi_1^0 \leftrightarrow u_{Rj}$, involving the neutral up-like higgs, or $d_{Li} + \phi_2^0 \leftrightarrow d_{Rj}$, involving the neutral down-like higgs, ensures that only one chemical potential, say that for u_L, is required to describe the quark sector. Similarly, efficient flavour mixing among the light neutrinos holds for mixing angles, $\theta_{ij} \gtrsim 10^{-7}$, so that only one chemical potential, μ_{ν_L}, is needed to describe the neutrino and sectors. In addition non-perturbative EW processes can be thought of creating, say, a neutron made of lhd quarks, and a lhd neutrino out of the vacuum for each generation. When this process is in equilibrium it imposes an additional relation which can be written

$$3\mu_{u_L} + \mu_{\nu_L} = 0. \tag{11}$$

When the MSSM susy spectrum is included the analysis becomes more involved. To simplify the analysis let us first assume at a temperature T (with $T > T_c$) that the MSSM susy spectrum, including k rhd sneutrinos can be considered light ($m \lesssim T$). As shown in Ref.[16] the usual soft-susy-breaking terms and supersymmetric partners of the SM interactions lead to enough relations among the chemical potentials that only 3 are needed for an independent set: $\mu_0 \equiv (\mu_{\phi_{10}} - \mu_{\phi_{20}})/2$ in the neutral Higgs sector, μ_{u_L}, and μ_{ν_L}. In our case we must add an additional k light ($m < T_c$) rhd neutrinos and sneutrinos with their associated potentials $\mu_{\nu_{Ri}}$ and $\mu_{\tilde{\nu}_{Ri}}$, but we also have A-term interactions between the lhd and rhd sneutrinos, and the very small lepton-number violating rhd sneutrino B-term. These together enforce the conditions

$$\mu_{\tilde{\nu}_R} = \mu_{\tilde{\nu}_L} = \mu_{\nu_L}, \tag{12}$$

where the first equality arises from A-term mediated interactions (as the A terms are $\mathcal{O}(m_w)$ it is simple to show that these are in equilibrium), and the second equality arises from $SU(2)_w \times U(1)_Y$ gaugino interactions.[4] In terms of μ_0, μ_{u_L}, and μ_{ν_L}, the condition Eq.(11) is unchanged in the case of the MSSM $+ k$ rhd neutrino supermultiplets.

The final step in the calculation of the relative number densities proceeds by expressing the relevant densities in terms of the independent chemical potentials. For the baryon number density B the result is simply

$$B = 6T^2 \mu_{u_L}. \tag{13}$$

For the sneutrino asymmetry the result depends upon the mass spectrum of the sneutrino mass eigenstates relative to T_c. In the range of parameters of interest to us where the light sneutrino eigenstate $\tilde{\nu}_1$ is the LSP we have $m_1 < T_c$. On the other hand the heavier eigenstate $\tilde{\nu}_2$ can have $m_2 > T_c$ or $m_2 < T_c$. If both eigenstates satisfy $m < T_c$ then the sneutrino density is (allowing for a possible k such states)

$$S = \frac{kT^2}{3}(\mu_{\tilde{\nu}_L} + \mu_{\tilde{\nu}_R}) = \frac{2kT^2}{3}\mu_{\nu_L} \tag{14}$$

while if only $m_1 < T_c$ then S is half as large. Finally using the non-perturbative relation Eq.(11) we find that

$$\frac{A}{A_{\text{bary}}} = -\frac{S}{B} = \frac{k}{3} \ \ \text{to} \ \ \frac{k}{6} \tag{15}$$

depending on m_2. We view $k = 1$ as most likely, so we specialize to that case in the next section. Finally we mention that a full susy-mass-threshold dependent calculation of the relative asymmetries introduces only small corrections to this result (see Ref.[16] for a similar analysis in the MSSM case).

In the end the vital point is that independent of whatever dynamics at scales $E > T_c$ produces an asymmetry in either the baryon sector, or in the neutrino or sneutrino sector, the $(B+L)$-anomaly-induced interactions together with EW gaugino and A-term interactions automatically distributes the asymmetry between the baryons and the light *sneutrino* states, with a predictable A/A_{bary} ratio. The asymmetry could originate from, e.g. a GUT-based baryogenesis mechanism, or maybe more interestingly in the context of the sneutrino dark matter model there is the possibility that it could arise via a resonant leptogenesis mechanism at the TeV-scale as discussed in Ref.[10]. The end result is that we expect $1/3 \gtrsim A/A_{\text{bary}} \gtrsim 1/6$ independent of the source of the asymmetry.

4   Asymmetric sneutrino and baryon abundances

We will now apply the formalism of section 2 to the sneutrino model. Given the value of $1/3 \gtrsim A/A_{\text{bary}} \gtrsim 1/6$, we expect the role of the matter-antimatter asymmetry to be most

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[4]Note that the rhd neutrinos only interact via the small neutrino Yukawa, $\lambda \sim 10^{-7} - 10^{-8}$, so these states are not in equilibrium at $T \sim T_c$. However, μ_{ν_R} does not enter our final expression for the sneutrino to baryon number density ratio, so its value is irrelevant.

important for sneutrinos with masses on the order of $m \sim \text{few} \times m_p \, \Omega_{\text{WMAP}} h^2 / \Omega_{\text{bary}} h^2 \simeq$ 30 GeV. For this reason, we will focus on the mass range of about 10 to 80 GeV. In this range, the most important annihilation channel is s-channel Z-exchange to fermions. This is a p-wave amplitude, and has no a-term in the expansion in v^2 (see section 2). This is also only a sneutrino-antisneutrino process (*i.e.* $\sigma_{11,22}$, not $\sigma_{12,21}$). Other channels include t-channel neutralino and chargino exchanges to neutrinos and charged leptons. The channel to charged leptons includes both a and b terms in the v^2 expansion. Additionally, sneutrino-sneutrino pairs (or antisneutrino-antisneutrino pairs) can annihilate via t-channel neutralino exchange to neutrino pairs of the same flavor. This is the only contribution to $\sigma_{12,21}$ in this model.

The annihilation cross sections for the neutralino and chargino exchange processes depend on the masses and mixings of these exchanged particles, although the Z-exchange channels do not. We have adopted the parameters M_1=300 GeV, M_2=300 GeV, μ=600 GeV and $\tan \beta = 50$ throughout our calculation. Additionally, there is the mixing angle between the sneutrino and the singlet which can be varied. This quantity has the effect of reducing all electroweak couplings of the sneutrino eigenstate by a factor of $\sin^2 \theta$, *i.e.* reducing the annihilation cross sections by a factor of $\sin^4 \theta$.

Finally, one can consider the s-channel exchange of Higgs bosons to $b\bar{b}$ (or ZZ and W^+W^- above the appropriate mass thresholds). This process is not suppressed by $\sin^4 \theta$. The efficiency of this process depends on the magnitude of the trilinear scalar coupling and is only important near the Higgs resonance ($m_{\tilde{\nu}} \simeq m_h/2$) or above the gauge boson mass thresholds.

Our results are shown in figure 1. The shaded regions of the parameter space predict a relic density within the range measured by WMAP ($0.129 > \Omega_{\text{CDM}} h^2 > 0.095$). In the left frame, no matter-antimatter asymmetry was included. This frame can be easily understood as requiring greater suppression of electroweak couplings (smaller $\sin \theta$) as the sneutrino mass approaches the Z mass. The dip at 56-59 GeV is due to s-channel higgs exchange to $b\bar{b}$. In the right frame, a matter-antimatter asymmetry of $A/A_{\text{bary}} \simeq 1/6$ was included. To further illustrate this effect, we show the result of this calculation across one value of $\sin \theta$ in figure 2. Below about 30 GeV, the matter-antimatter asymmetry has little effect on the calculation and the solid and dot-dashed lines fall nearly on top of each other. In the range of roughly 30-70 GeV, however, the asymmetry pulls the relic density above the standard symmetric result into the range favored by WMAP. Above this range, sneutrino-antisneutrino annihilation decreases, leading to larger relic densities for the case with no asymmetry. The relic density for the asymmetric case, however, is largely determined by the sneutrino-sneutrino annihilation cross section in this region, so does not increase as rapidly, therefore resulting in a relic density much closer to the preffered value, even for $m_{\tilde{\nu}} > 70$ GeV.

5   Indirect Detection

Sneutrinos in the mass range we have studied here annihilate dominantly through s-channel Z exchange in the early universe. This annihilation channel consists of a b-term
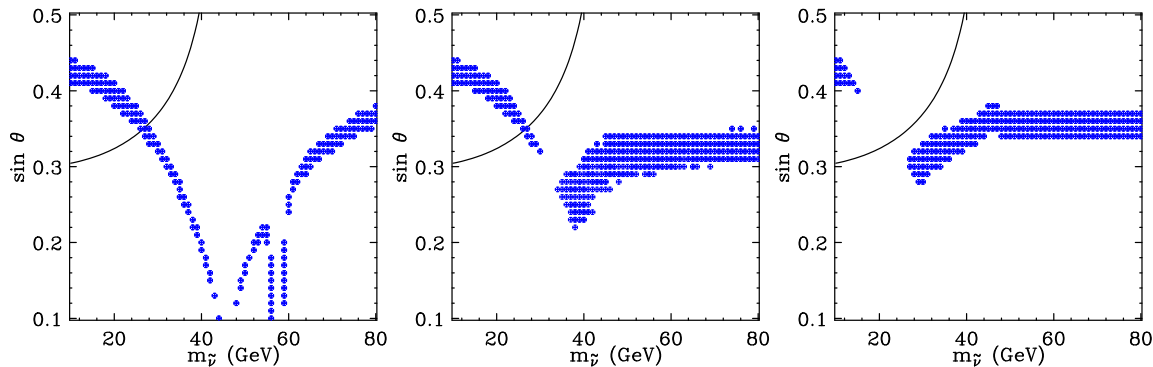
Figure 1: The regions of parameter space which provide the quantity of mixed sneutrino cold dark matter measured by WMAP, $0.129 > \Omega_{\mathrm{CDM}} h^2 > 0.095$. In the left frame, the standard calculation with no matter-antimatter asymmetry is used. In the center frame the effect of a matter-antimatter asymmetry with $A/A_{\mathrm{bary}} \simeq 1/6$ is included, while in the right frame a matter-antimatter asymmetry $A/A_{\mathrm{bary}} \simeq 1/3$ is assumed. Notice the much larger regions of acceptable parameter space in cases with asymmetry. We use the following parameters: M_1=300 GeV, M_2=300 GeV, μ=600 GeV, $\tan\beta = 50$ and m_h=115 GeV. The region above the solid line is excluded by measurements of the invisible Z decay width [14].

in the expansion $\sigma v = a + b v^2 + \mathcal{O}(v^4)$, and thus is very efficient for annihilations at low velocities which take place in the late universe. At low velocities, the process $\tilde{\nu}\tilde{\nu} \to \nu\nu$, via t-channel neutralino exchange, is the most important annihilation channel.

Given these characteristics, searches for gamma-rays, anti-protons, positrons or other particles produced in sneutrino annihilations in the Galactic halo are quite inefficient. Due to the rather large elastic scattering cross sections of sneutrinos, however, indirect detection techniques involving the capture of dark matter in bodies such as the Sun can be very effective.

Neutrino telescopes detect neutrinos produced in dark matter annihilations in the core of the Sun by observing 'tracks' of muons generated in charged current interactions. If our sneutrino dark matter candidate is of muon or tau flavor, the flux of muons generated in this way in a detector at Earth is on the order of 10^{-12} cm^{-2} s^{-1} [8], in conflict with data from the Super-Kamiokande and MACRO experiments, for example [17]. If the sneutrino candidate, however, is of electron flavor, these bounds can be easily averted.

6   Conclusions

In the standard freeze-out calculation for a weakly interacting dark matter relic, there is little reason to expect a density of dark matter which is similar to the density of baryons, short of anthropic arguments. In this letter, however, we have presented an explanation for the similarity between these two quantities. Our solution introduces an asymmetry between dark matter particles and anti-particles which is related to the
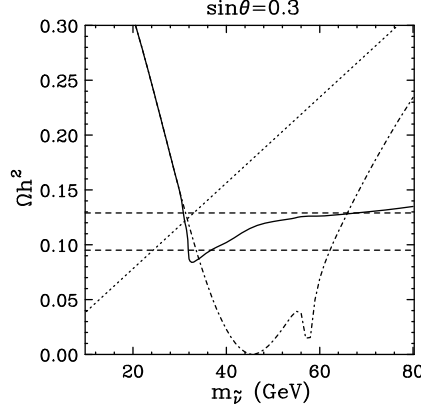
Figure 2: The thermal relic density as a function of mass for sneutrinos+antisneutrinos with no asymmetry (dot-dash), with a matter-antimatter asymmetry of $A/A_{\mathrm{bary}} \simeq 1/6$ (solid) and the estimate of Eq.(5) (dots). The relic density range favored by WMAP is bound by dashed lines ($0.129 > \Omega_{\mathrm{CDM}} h^2 > 0.095$). $\sin \theta = 0.3$, M_1=300 GeV, M_2=300 GeV, μ=600 GeV, $\tan \beta = 50$ and m_h=115 GeV have been used. See text for details.

baryon-antibaryon asymmetry. This leads to a natural dark matter relic density of the same order of magnitude as the baryon density.

As an example, we considered a mixed sneutrino dark matter candidate which transfers its particle-antiparticle asymmetry to the baryons through the electroweak anomaly. We carry out the relic density calculation for such a candidate and find substantial regions of parameter space in which the observations of WMAP can be satisfied. With no asymmetry, only a narrow strip of parameter space can provide the observed relic density.

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Asymmetric Sneutrino Dark Matter and the $\Omega_{\mathrm{b}}/\Omega_{\mathrm{DM}}$ Puzzle

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Abstract

The inferred values of the cosmological baryon and dark matter densities are strikingly similar, but in most theories of the early universe there is no true explanation of this fact; in particular, the baryon asymmetry and thus density depends upon unknown, and *a priori* unknown and possibly small, CP-violating phases which are independent of all parameters determining the dark matter density. We consider models of dark matter possessing a particle-antiparticle asymmetry where this asymmetry determines both the baryon asymmetry and strongly effects the dark matter density, thus naturally linking Ω_{b} and Ω_{dm}. We show that sneutrinos can play the role of such dark matter in a previously studied variant of the MSSM in which the light neutrino masses result from higher-dimensional supersymmetry-breaking terms.

1 Introduction

For some time it has been apparent that the inferred values of the cosmological baryon and dark matter densities are strikingly similar. For example, the latest WMAP-determined range for the dark matter density, $0.129 > \Omega_{\mathrm{CDM}}h^2 > 0.095$, is within a factor of a few of the combined WMAP and big-bang nucleosynthesis determined value of the baryon density, $0.025 > \Omega_{\mathrm{b}}h^2 > 0.012$ [1, 2]. In the vast majority of models of the early universe, the cosmological baryon and dark matter densities are independently determined. The surviving baryon density is set by a baryon asymmetry generated during baryogenesis, and thus depends upon unknown, and *a priori* unknown and possibly small CP-violating phases, as well as unknown baryon-number violating dynamics. In contrast, the dark matter density is set by the 'freeze-out' of the interactions that keep the dark matter in equilibrium, and is independent of the dynamics of baryogenesis. In particular, although weakly-interacting massive particles at the TeV-scale naturally have a relic density of $\mathcal{O}(1)$ times the critical density, this is not at all the case for baryons. The Boltzmann equations show the size of the baryon relic density in the absence of an asymmetry is $\mathcal{O}(10^{-11})$. Thus the comparative closeness of the baryon and dark matter densities poses a puzzle. One possible approach to this problem is to more closely integrate the dynamics of baryogenesis with that of the origin of dark matter. In particular, it is natural to consider models where the dark matter and baryon sectors share a quantum number, either continuous or discrete, which provides a relation between their surviving number densities and thus energy densities.[1]

Specifically, in this letter we consider models of dark matter possessing a particle-antiparticle asymmetry where this asymmetry strongly effects the dark matter density, and through the electroweak (EW) anomaly, determines the baryon asymmetry, thus naturally linking Ω_{b} and Ω_{CDM}. (A noteworthy early attempt along these lines which shares some features with our model is contained in Ref. [9].) The shared quantum number clearly requires that the dark matter particle not be it's own antiparticle, so the putative dark matter candidate cannot be the standard LSP neutralino of the broken Minimal Supersymmetric Standard Model (MSSM). In fact, we will show in sections 3 and 4 that sneutrinos can play the role of such dark matter in a previously studied variant of the MSSM in which the light neutrino masses result from higher-dimensional supersymmetry-breaking terms [10, 11, 12, 13, 14]. This model preserves all the successes of the MSSM, such as stability of the weak scale and unification of gauge couplings, while being, at least in part, testable at the LHC.

Before we proceed to the calculation of the relic dark matter and baryon densities and the details of our model, it is useful to consider some aspects of the idea of a shared quantum number determining the ratio of baryon to dark matter. First, consider the simplified case in which, to a very good approximation, the dark matter and baryon sectors cannot exchange their conserved quantum number after its first production. As an example, consider the situation in which the asymptotic baryon number $B = 1$ states of the Standard Model (SM) have global charge, q, while the lightest asymptotic states

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[1]An alternate approach is to invoke a form of the Anthropic Principle, see, for example, Ref. [3]. Other ideas are explored in Refs. [4, 5, 6, 7, 8].

(with mass m_{dm}) in a hidden sector carry charge, Q. Then conservation of global charge implies that $q\,(n_{\mathrm{b}} - n_{\overline{\mathrm{b}}}) = -Q\,(n_{\mathrm{dm}} - n_{\overline{\mathrm{dm}}})$, where the n's are the number densities of the indicated (anti)particles. Further assume that interactions in both sectors are strong enough such that almost all antiparticles are eliminated by annihilation with their particles (assuming $Q/q < 0$ for simplicity). This implies

$$n_{\mathrm{b}} = c\, n_{\mathrm{dm}} \quad \text{with} \quad c = |Q/q|, \tag{1}$$

which in turn leads to

$$\frac{\Omega_{\mathrm{b}}}{\Omega_{\mathrm{dm}}} = \frac{m_{\mathrm{b}}}{m_{\mathrm{dm}}} \frac{n_{\mathrm{b}}}{n_{\mathrm{dm}}} = c\,\frac{m_{\mathrm{b}}}{m_{\mathrm{dm}}}. \tag{2}$$

Thus the energy densities are indeed related, but the ratio is only naturally $\mathcal{O}(1)$ if the ratio of the baryon to dark matter masses is not too small. It is important to note that this disfavours models of the above type where the dark matter candidate is a 'hidden sector' particle and favours a particular class of models where the dark matter candidate has weak scale mass, its mass arising either from electroweak symmetry breaking, or from the dynamics which drives electroweak breaking, such as supersymmetry softly broken at the weak-scale.[2]

As a consequence, the naive statement of Eq.(1) also requires modification as the approximation of negligible interactions between the two sectors, and thus negligible exchange of the shared conserved quantum number, does not hold in realistic models. In this situation, the ratio of the conserved quantum number stored in the two sectors, c, is not simply given by Q/q, but is determined by the 'chemical' equilibrium conditions between the two sectors just before the freeze-out of the relevant interactions. We will further discuss this point in section 2. Moreover, two other alterations to the naive relation Eq.(2) can be present. First, the shared quantum number may only be conserved modulo n, in which case 'self' annihilations may occur, and second, it is not always the case that dark matter interactions are efficient enough such that $n_{\overline{\mathrm{dm}}} \ll n_{\mathrm{dm}}$. Nevertheless, as we argue in the next section, a careful consideration of such models shows that a version of the naive statement Eq.(2) can indeed hold.

2 The Effect of a Matter-Antimatter Asymmetry on the Density of a Thermal Relic

If there were an asymmetry between the density of dark matter particles and dark antimatter particles in the early universe, it may have a considerable effect on the relic density calculation [15]. In this section, we review the relic density calculation for a single species, two species with no asymmetry and two species with an asymmetry.

If dark matter consists of a single species, the relic density calculation is straightforward. The most common example of a single species dark matter candidate would be

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[2]In extensions of the SM such as the MSSM where many gauge-non-neutral fermionic and scalar states also gain mass at this scale, the beta-function coefficient for $SU(3)_{\mathrm{QCD}}$ sharply increases as energies are reduced below the weak scale, and the confinement scale and thus baryon mass is naturally not very much smaller than m_{weak}.

a neutralino LSP with negligible coannihilations with other SUSY particles. In such a case, the resulting relic density is simply

$$\Omega h^2 \cong \frac{1.04 \times 10^9}{M_{\mathrm{Pl}}} \frac{x_F}{\sqrt{g_*}} \frac{1}{(a + 3b/x_F)}, \tag{3}$$

where M_{Pl} is the Planck mass, x_F is the WIMP's mass over the freeze-out temperature, g_* is number of relativistic degrees of freedom available at freeze-out and a and b are related to the WIMP's self-annihilation cross section by $\sigma v = a + bv^2 + \mathcal{O}(v^4)$, v being the relative velocity of the annihilating WIMPs. For a weakly interacting relic, x_F is generally about 20 and varies only mildly with changes in the cross section or mass.

If we instead consider a dark matter candidate *and* its antiparticle, the calculation becomes somewhat different. Labeling these species 1 and 2, we now have each particle's self-annihilation cross section, $\sigma_{11} = \sigma_{22}$, and the particle-antiparticle annihilation cross section, $\sigma_{12} = \sigma_{21}$. In the limit that σ_{12} is much smaller than σ_{11}, we simply get two decoupled species which each provide the relic density of Eq.(3). Otherwise, the effective cross section becomes $\sigma_{11} + \sigma_{12}$ for each species. For the two species together, a density of

$$\Omega h^2 \cong 2 \times \frac{1.04 \times 10^9}{M_{\mathrm{Pl}}} \frac{x_F}{\sqrt{g_*}} \frac{1}{(a_{11} + a_{12} + 3(b_{11} + b_{12})/x_F)} \tag{4}$$

is produced. This case is essentially the calculation of the density of a relic with the coannihilation of a degenerate particle [16].

If an asymmetry exists between the number density of particles and antiparticles in the early universe, again this calculation becomes modified. The most well known example of this is the baryon-antibaryon asymmetry which leads to the current baryon density of the universe despite the very large baryon-antibaryon annihilation cross section. In such a case, where the particle-particle and antiparticle-antiparticle annihilation cross sections (σ_{11} and σ_{22}) are zero, or negligibly small, the minimum relic density can be related to the magnitude of the particle-antiparticle asymmetry:

$$\Omega h^2_{\mathrm{min}} = \Omega h^2_{\mathrm{bary}} \frac{A}{A_{\mathrm{bary}}} \frac{m}{m_{\mathrm{bary}}}, \tag{5}$$

where $\Omega h^2_{\mathrm{bary}}$ is the current density of baryons in the universe, A and A_{bary} are the particle-antiparticle asymmetries of our relic and of baryons, defined by $A = (n - \overline{n})/n$. A_{bary} is known to be of the order of 10^{-9}. m and m_{bary} are the masses of our relic and of baryons (*i.e.* the proton mass). Of course, if the particle-particle annihilation cross section for the relic is not negligible, Eq.(5) will not hold. There will be a generic tendency for the density of a relic to move towards this value as a result of an asymmetry, however.

To accurately assess the effect of an asymmetry on the thermal relic abundance, the effective cross sections must be determined by integrating over the thermal history. In particular, for each species separately, the densities predicted are

$$\Omega_1 h^2 \cong \frac{1.04 \times 10^9}{M_{\mathrm{Pl}}} \frac{x_F}{\sqrt{g_*}} \frac{1}{(a_{11,\mathrm{eff}} + a_{12,\mathrm{eff}} + 3(b_{11,\mathrm{eff}} + b_{12,\mathrm{eff}})/x_F)} \tag{6}$$

and

$$\Omega_2 h^2 \cong \frac{1.04 \times 10^9}{M_{\text{Pl}}} \frac{x_F}{\sqrt{g_*}} \frac{1}{(a_{22,\text{eff}} + a_{21,\text{eff}} + 3(b_{22,\text{eff}} + b_{21,\text{eff}})/x_F)}, \qquad (7)$$

where the effective cross sections are given by

$$
\begin{array}{llll}
a_{11,\text{eff}} & = & a_{11}, & a_{12,\text{eff}} = x_F \int_{x_F}^{\infty} \frac{a_{12}}{x^2} \left(\frac{e^{-x}}{e^{-x}+A} \right) dx, \\
a_{22,\text{eff}} & = & a_{22}, & a_{21,\text{eff}} = x_F \int_{x_F}^{\infty} \frac{a_{12}}{x^2} \left(\frac{e^{-x}+A}{e^{-x}} \right) dx, \\
b_{11,\text{eff}} & = & b_{11}, & b_{12,\text{eff}} = 2x_F^2 \int_{x_F}^{\infty} \frac{b_{12}}{x^3} \left(\frac{e^{-x}}{e^{-x}+A} \right) dx, \\
b_{22,\text{eff}} & = & b_{22}, & b_{21,\text{eff}} = 2x_F^2 \int_{x_F}^{\infty} \frac{b_{12}}{x^3} \left(\frac{e^{-x}+A}{e^{-x}} \right) dx.
\end{array}
\qquad (8)
$$

The factor, $e^{-x}/(e^{-x} + A)$, is the depletion of the second species due to the asymmetry at a temperature, $T = m/x$. The effect of this depletion depends on the size of the asymmetry and the cross sections involved. If the estimate of Eq.(5) is considerably less than the standard two-species result of Eq.(4), then the asymmetry has little impact and the sum of Eqs.(6) and (7) is nearly equal to Eq.(4). If, however, the estimate of Eq.(5) is larger than the standard two-species result, the relic density moves from the value found with Eq.(4) towards that of Eq.(5). If σ_{11}/σ_{12} is small, then the result will be closer to Eq.(5). If σ_{11}/σ_{12} is large, then the density will be less modified. Note that in some cases, the effect of an asymmetry can be to lower the density of a relic by up to a factor of two. This occurs when σ_{11} and σ_{12} are comparable thus efficiently deplete the second species while leaving the first largely uneffected.

Although these effects are rather model dependent, the net effect can be large regions of parameter space which predict the observed density of dark matter for a WIMP with a mass near $m \sim m_p \Omega_{\text{WMAP}} h^2 / \Omega_{\text{bary}} h^2$, for a matter-antimatter asymmetry similar to that for baryons. For concreteness, we consider a specific example of such a model in the following section.

3   An Example: Mixed Sneutrino Dark Matter

Within the context of the Minimal Supersymmetric Standard Model (MSSM), sneutrinos do not make a very appealing dark matter candidate. They annihilate quite efficiently, resulting in a relic density smaller than the observed dark matter density for sneutrinos lighter than about 500 GeV. Furthermore, their elastic scattering cross section is sufficiently large to be easily observed by direct dark matter experiments. Thus it is natural that in contrast to other supersymmetric dark matter candidates, sneutrino dark matter has received considerably less attention [17].

Looking beyond the vanilla MSSM, however, it is quite easy to construct a viable model for sneutrino dark matter. For example, sneutrinos could mix through A-terms with a singlet forming a mixed sneutrino state. This would reduce the annihilation cross section, potentially providing the appropriate quantity of dark matter, as well as reduce the elastic scattering cross section to experimentally acceptable levels. In fact, such a

scenario has previously been discussed as a way to reconcile the results of the direct dark matter detection experiments DAMA and CDMS [12]. These models have light sneutrino mass eigenstates which are mixtures of left-handed 'active' sneutrinos, and right-handed 'sterile' sneutrinos.

Such mixing naturally arises in models in which the light neutrino masses are generated by supersymmetry breaking effects [10, 11, 13, 14]. The importance to us of these models is that the light sneutrino states share a non-anomalous $(B-L)$-symmetry with the baryons which is only weakly broken by the Majorana neutrino masses. In analogy to the Giudice-Masiero mechanism for the Higgs μ-term there exists a global symmetry that forbids a direct mass, $M_R NN$, and neutrino Yukawa coupling, λLNH_u, involving the SM singlet right-handed neutrino superfield N, but allows non-renormalizable couplings to a supersymmetry-breaking hidden-sector field T. The right-handed neutrino mass and the neutrino Yukawa coupling, λLNH_u, then only arise from $1/M_{\mathrm{planck}}$-suppressed terms involving the fundamental supersymmetry breaking scale $m_I \sim (vM_{pl})^{1/2}$, where v is the weak scale. The relevant terms in the Lagrangian are

$$\mathcal{L} = \int d^4\theta \left(h\frac{T^\dagger}{M}NN + \tilde{h}\frac{T^\dagger T}{M^2}N^\dagger N + h_B \frac{T^\dagger T T^\dagger}{M^3}NN + \dots \right) + \int d^2\theta \left(g\frac{T}{M}LNH_u \right), \quad (9)$$

where all dimensionless couplings g, h, etc, are taken to be $\mathcal{O}(1)$. The structure of the Lagrangian can be enforced by some simple symmetries [10, 11, 13], including R-parity, which leads to a stable lightest supersymmetric particle (LSP) dark matter candidate. After supersymmetry is broken in the hidden sector at the scale m_I, the F and A components of the field T acquire intermediate scale vevs, $\langle F_T \rangle \sim \langle A_T \rangle^2 \sim m_I^2$. Substituting these vevs into the Lagrangian shows that after SUSY breaking, there arises: (1) a neutrino Yukawa coupling of size $\sim 10^{-7} - 10^{-8}$, (2) right-handed neutrino and sneutrino masses at the *weak scale*, (3) a weak-scale trilinear scalar A-term, and (4) a right-handed sneutrino lepton-number violating B-term with magnitude $B^2 \sim (\mathrm{few} \times 100\,\mathrm{MeV})^2$. These together lead to two sources of neutrino masses, a tree level see-saw contribution, as well as a dominant 1-loop contribution, leading to a consistent and attractive neutrino phenomenology [10, 13].

More relevant for us, the sneutrino mass matrix squared is

$$\left(\begin{array}{cc} \tilde{\nu}^* & n \end{array} \right) \left(\begin{array}{cc} m_L^2 & \frac{1}{\sqrt{2}}Av\sin\beta \\ \frac{1}{\sqrt{2}}Av\sin\beta & m_R^2 \end{array} \right) \left(\begin{array}{c} \tilde{\nu} \\ n^* \end{array} \right), \quad (10)$$

where we have ignored the small lepton number violating B-term for the moment. Therefore mixing between the left-handed 'active' sneutrino, $\tilde{\nu}$, and the right-handed 'sterile' sneutrino, $\tilde{n}$, occurs, producing the light mass eigenstate $\tilde{\nu}_1 = -\tilde{\nu}\sin\theta + \tilde{n}^*\cos\theta$, where $\tan 2\theta = -\sqrt{2}Av\sin\beta/(m_R^2 - m_L^2)$, with $m_R^2 = M_N^2 + m_{R,\mathrm{soft}}^2$ and $m_L^2 = m_Z^2\cos 2\beta/2 + m_{L,\mathrm{soft}}^2$. Since the coupling of the lighter sneutrino eigenstate to the Z is suppressed by $\sin\theta$, the direct LEP experimental constraints are weakened, and the calculation of the sneutrino relic abundances is modified. A numerical analysis of the susy parameter space shows that a mixing $\sin\theta \simeq 0.3$ is quite natural.[3] (The B^2 mass-term splits the CP-even and CP-odd components of $\tilde{\nu}_1$ by the amount $\delta \sim \cos^2\theta B^2/m_1$.

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[3]Note that the size of the A terms is limited by considerations of vacuum stability [14].

This can have important consequences for detection, potentially making consistent the DAMA and CDMS experiments depending upon parameters [12]. For the purposes of our discussion this splitting is irrelevant.)

We now turn to the calculation of the relative asymmetry in the sneutrino and baryon sectors. The method is a simple adaption of the standard 'chemical' equilibriation techniques applied to, for example, the calculation of the ratio $B/(B-L)$ in the SM [19] or MSSM [20] in the presence of anomaly-induced baryon number violating processes in the early universe. For each species of particle we introduce a chemical potential, μ_A, and for each interaction in equilibrium, say $A + B \leftrightarrow C + D$ there is a corresponding relation among chemical potentials $\mu_A + \mu_B = \mu_C + \mu_D$. Note that when in *thermal* (not necessarily chemical) equilibrium the CP-conjugate species $\overline{A}$ has $\mu_{\overline{A}} = -\mu_A$. The asymmetry in particle-antiparticle number density for the relevant case of $\mu \ll T$ and a species of mass m is given by $n - \overline{n} = gT^2\mu/6$ for $T \gg m$, and is Boltzmann suppressed by $\exp(-m/T)$ for $T \lesssim m$.

The electroweak $B+L$ anomaly interactions provide an efficient mechanism to transmute between B- and L-number densities, at least above the EW phase transition temperature T_c where they are rapid. Our task is to calculate the ratio of sneutrino to baryon asymmetries at T_c where the exchange between the two sectors switches off. (We implicitly assume that the EWPT is 1st order so that anomaly-induced interactions drop out of equilibrium below T_c. For realistic values of T_c the results of our calculation are not particularly sensitive to this assumption.) Above the EWPT temperature the chemical potentials for all SM gauge bosons vanish, so, for example, members of an EW doublet have the same μ. Moreover, efficient flavour mixing among quarks induced by EW interactions such as $u_{Li} + \phi_1^0 \leftrightarrow u_{Rj}$, involving the neutral up-like higgs, or $d_{Li} + \phi_2^0 \leftrightarrow d_{Rj}$, involving the neutral down-like higgs, ensures that only one chemical potential, say that for u_L, is required to describe the quark sector. Similarly, efficient flavour mixing among the light neutrinos holds for mixing angles, $\theta_{ij} \gtrsim 10^{-7}$, so that only one chemical potential, μ_{ν_L}, is needed to describe the neutrino and sectors. In addition non-perturbative EW processes can be thought of creating, say, a neutron made of lhd quarks, and a lhd neutrino out of the vacuum for each generation. When this process is in equilibrium it imposes an additional relation which can be written

$$3\mu_{u_L} + \mu_{\nu_L} = 0. \tag{11}$$

When the MSSM susy spectrum is included the analysis becomes more involved. To simplify the analysis let us first assume at a temperature T (with $T > T_c$) that the MSSM susy spectrum, including k rhd sneutrinos can be considered light ($m \lesssim T$). As shown in Ref.[20] the usual soft-susy-breaking terms and supersymmetric partners of the SM interactions lead to enough relations among the chemical potentials that only 3 are needed for an independent set: $\mu_0 \equiv (\mu_{\phi_{10}} - \mu_{\phi_{20}})/2$ in the neutral Higgs sector, μ_{u_L}, and μ_{ν_L}. In our case we must add an additional k light ($m < T_c$) rhd neutrinos and sneutrinos with their associated potentials $\mu_{\nu_{Ri}}$ and $\mu_{\tilde{\nu}_{Ri}}$, but we also have A-term interactions between the lhd and rhd sneutrinos, and the very small lepton-number violating rhd sneutrino B-term. These together enforce the conditions

$$\mu_{\tilde{\nu}_R} = \mu_{\tilde{\nu}_L} = \mu_{\nu_L}, \tag{12}$$

6

where the first equality arises from A-term mediated interactions (as the A terms are $\mathcal{O}(m_w)$ it is simple to show that these are in equilibrium), and the second equality arises from $SU(2)_w \times U(1)_Y$ gaugino interactions.[4] In terms of μ_0, μ_{u_L}, and μ_{ν_L}, the condition Eq.(11) is unchanged in the case of the MSSM $+ k$ rhd neutrino supermultiplets.

The final step in the calculation of the relative number densities proceeds by expressing the relevant densities in terms of the independent chemical potentials. For the baryon number density B the result is simply

$$B = 6T^2 \mu_{u_L}. \tag{13}$$

For the sneutrino asymmetry the result depends upon the mass spectrum of the sneutrino mass eigenstates relative to T_c. In the range of parameters of interest to us where the light sneutrino eigenstate $\tilde{\nu}_1$ is the LSP we have $m_1 < T_c$. On the other hand the heavier eigenstate $\tilde{\nu}_2$ can have $m_2 > T_c$ or $m_2 < T_c$. If both eigenstates satisfy $m < T_c$ then the sneutrino density is (allowing for a possible k such states)

$$S = \frac{kT^2}{3}(\mu_{\tilde{\nu}_L} + \mu_{\tilde{\nu}_R}) = \frac{2kT^2}{3}\mu_{\nu_L} \tag{14}$$

while if only $m_1 < T_c$ then S is half as large. Finally using the non-perturbative relation Eq.(11) we find that

$$\frac{A}{A_{\text{bary}}} = -\frac{S}{B} = \frac{k}{3} \quad \text{to} \quad \frac{k}{6} \tag{15}$$

depending on m_2. We view $k = 1$ as most likely, so we specialize to that case in the next section. Finally we mention that a full susy-mass-threshold dependent calculation of the relative asymmetries introduces only small corrections to this result (see Ref.[20] for a similar analysis in the MSSM case).

In the end the vital point is that independent of whatever dynamics at scales $E > T_c$ produces an asymmetry in either the baryon sector, or in the neutrino or sneutrino sector, the $(B+L)$-anomaly-induced interactions together with EW gaugino and A-term interactions automatically distributes the asymmetry between the baryons and the light *sneutrino* states, with a predictable A/A_{bary} ratio. The asymmetry could originate from, e.g. a GUT-based baryogenesis mechanism, or maybe more interestingly in the context of the sneutrino dark matter model there is the possibility that it could arise via a resonant leptogenesis mechanism at the TeV-scale as discussed in Ref.[14]. The end result is that we expect $1/3 \gtrsim A/A_{\text{bary}} \gtrsim 1/6$ independent of the source of the asymmetry.

4 Asymmetric sneutrino and baryon abundances

We will now apply the formalism of section 2 to the sneutrino model. Given the value of $1/3 \gtrsim A/A_{\text{bary}} \gtrsim 1/6$, we expect the role of the matter-antimatter asymmetry to be most

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[4]Note that the rhd neutrinos only interact via the small neutrino Yukawa, $\lambda \sim 10^{-7} - 10^{-8}$, so these states are not in equilibrium at $T \sim T_c$. However, μ_{ν_R} does not enter our final expression for the sneutrino to baryon number density ratio, so its value is irrelevant.

important for sneutrinos with masses on the order of $m \sim \text{few} \times m_p\, \Omega_{\text{WMAP}} h^2 / \Omega_{\text{bary}} h^2 \simeq$ 30 GeV. For this reason, we will focus on the mass range of about 10 to 80 GeV. In this range, the most important annihilation channel is s-channel Z-exchange to fermions. This is a p-wave amplitude, and has no a-term in the expansion in v^2 (see section 2). This is also only a sneutrino-antisneutrino process (*i.e.* $\sigma_{11,22}$, not $\sigma_{12,21}$). Other channels include t-channel neutralino and chargino exchanges to neutrinos and charged leptons. The channel to charged leptons includes both a and b terms in the v^2 expansion. Additionally, sneutrino-sneutrino pairs (or antisneutrino-antisneutrino pairs) can annihilate via t-channel neutralino exchange to neutrino pairs of the same flavor. This is the only contribution to $\sigma_{12,21}$ in this model.

The annihilation cross sections for the neutralino and chargino exchange processes depend on the masses and mixings of these exchanged particles, although the Z-exchange channels do not. We have adopted the parameters M_1=300 GeV, M_2=300 GeV, μ=600 GeV and $\tan\beta = 50$ throughout our calculation. Additionally, there is the mixing angle between the sneutrino and the singlet which can be varied. This quantity has the effect of reducing all electroweak couplings of the sneutrino eigenstate by a factor of $\sin^2\theta$, *i.e.* reducing the annihilation cross sections by a factor of $\sin^4\theta$.

Finally, one can consider the s-channel exchange of Higgs bosons to $b\bar{b}$ (or ZZ and W^+W^- above the appropriate mass thresholds). This process is not suppressed by $\sin^4\theta$. The efficiency of this process depends on the magnitude of the trilinear scalar coupling and is only important near the Higgs resonance ($m_{\tilde{\nu}} \simeq m_h/2$) or above the gauge boson mass thresholds.

Our results are shown in figure 1. The shaded regions of the parameter space predict a relic density within the range measured by WMAP ($0.129 > \Omega_{\text{CDM}} h^2 > 0.095$). In the left frame, no matter-antimatter asymmetry was included. This frame can be easily understood as requiring greater suppression of electroweak couplings (smaller $\sin\theta$) as the sneutrino mass approaches the Z mass. The dip at 56-59 GeV is due to s-channel higgs exchange to $b\bar{b}$. In the right frame, a matter-antimatter asymmetry of $A/A_{\text{bary}} \simeq 1/6$ was included. To further illustrate this effect, we show the result of this calculation across one value of $\sin\theta$ in figure 2. Below about 30 GeV, the matter-antimatter asymmetry has little effect on the calculation and the solid and dot-dashed lines fall nearly on top of each other. In the range of roughly 30-70 GeV, however, the asymmetry pulls the relic density above the standard symmetric result into the range favored by WMAP. Above this range, sneutrino-antisneutrino annihilation decreases, leading to larger relic densities for the case with no asymmetry. The relic density for the asymmetric case, however, is largely determined by the sneutrino-sneutrino annihilation cross section in this region, so does not increase as rapidly, therefore resulting in a relic density much closer to the preferred value, even for $m_{\tilde{\nu}} > 70$ GeV.

5   Indirect Detection

Sneutrinos in the mass range we have studied here annihilate dominantly through s-channel Z exchange in the early universe. This annihilation channel consists of a b-term
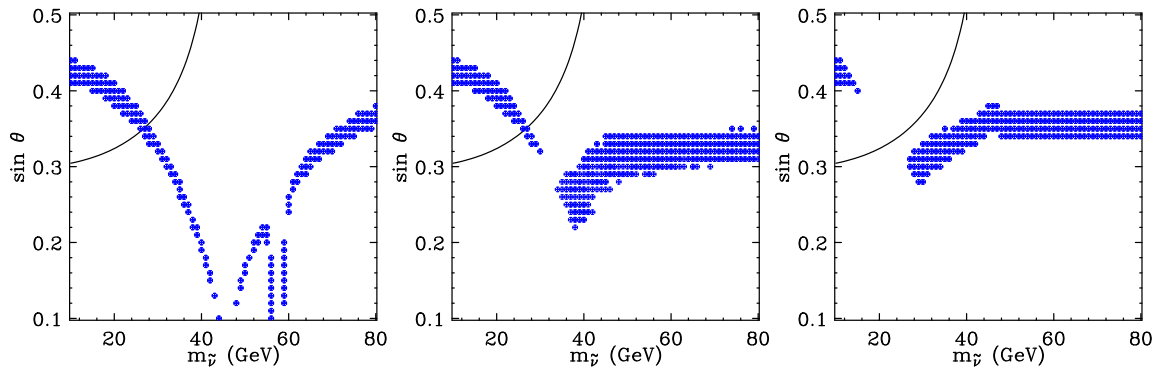
Figure 1: The regions of parameter space which provide the quantity of mixed sneutrino cold dark matter measured by WMAP, $0.129 > \Omega_{\mathrm{CDM}}h^2 > 0.095$. In the left frame, the standard calculation with no matter-antimatter asymmetry is used. In the center frame the effect of a matter-antimatter asymmetry with $A/A_{\mathrm{bary}} \simeq 1/6$ is included, while in the right frame a matter-antimatter asymmetry $A/A_{\mathrm{bary}} \simeq 1/3$ is assumed. Notice the much larger regions of acceptable parameter space in cases with asymmetry. We use the following parameters: M_1=300 GeV, M_2=300 GeV, μ=600 GeV, $\tan\beta = 50$ and m_h=115 GeV. The region above the solid line is excluded by measurements of the invisible Z decay width [18].

in the expansion $\sigma v = a + bv^2 + \mathcal{O}(v^4)$, and thus is very inefficient for annihilations at low velocities which take place in the late universe. At low velocities, the process $\tilde{\nu}\tilde{\nu} \to \nu\nu$, via t-channel neutralino exchange, is the most important annihilation channel.

Given these characteristics, searches for gamma-rays, anti-protons, positrons or other particles produced in sneutrino annihilations in the Galactic halo are quite inefficient. Due to the rather large elastic scattering cross sections of sneutrinos, however, indirect detection techniques involving the capture of dark matter in bodies such as the Sun can be very effective.

Neutrino telescopes detect neutrinos produced in dark matter annihilations in the core of the Sun by observing 'tracks' of muons generated in charged current interactions. If our sneutirno [5] dark matter candidate is of muon or tau flavor, the flux of muons generated in this way in a detector at Earth is of the order of 10^{-12} cm^{-2} s^{-1} [12], in conflict with data from the Super-Kamiokande and MACRO experiments, for example [21]. If the candidate, however, is of electron flavor, these bounds can be easily averted.

6   Conclusions

In the standard freeze-out calculation for a weakly interacting dark matter relic, there is little reason to expect a density of dark matter which is similar to the density of baryons, short of anthropic arguments. In this letter, however, we have presented an

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[5]In the model presented, the dark matter candidate is actually an antisneutrino and consequently antineutrinos will be produced by dark matter annihilations in the sun.
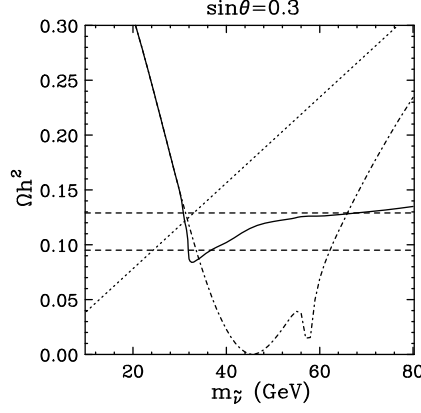
Figure 2: The thermal relic density as a function of mass for sneutrinos+antisneutrinos with no asymmetry (dot-dash), with a matter-antimatter asymmetry of $A/A_{\text{bary}} \simeq 1/6$ (solid) and the estimate of Eq.(5) (dots). The relic density range favored by WMAP is bound by dashed lines ($0.129 > \Omega_{\text{CDM}}h^2 > 0.095$). $\sin\theta$=0.3, M_1=300 GeV, M_2=300 GeV, μ=600 GeV, $\tan\beta = 50$ and m_h=115 GeV have been used. See text for details.

explanation for the similarity between these two quantities. Our solution introduces an asymmetry between dark matter particles and anti-particles which is related to the baryon-antibaryon asymmetry. This leads to a natural dark matter relic density of the same order of magnitude as the baryon density.

As an example, we considered a mixed sneutrino dark matter candidate which transfers its particle-antiparticle asymmetry to the baryons through the electroweak anomaly. We carry out the relic density calculation for such a candidate and find substantial regions of parameter space in which the observations of WMAP can be satisfied. With no asymmetry, only a narrow strip of parameter space can provide the observed relic density.

Acknowledgments

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